Note on Strange Quarks in the Nucleon

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ABSTRACT

Scalar matrix elements involving strange quarks are studied in several models. Apart from a critical reexamination of results obtained in the Nambu and Jona-Lasinio model we study a scenario, motivated by instanton physics, where spontaneous chiral symmetry breaking is induced by the flavor-mixing 't Hooft interaction only. We also investigate possible contributions of virtual kaon loops to the strangeness content of the nucleon (University of Regensburg preprint TPR-94-04).

Recently much interest has been focused on the role of strange quarks in the nucleon. The observables under discussion are the strange quark spin content of the proton [1, 2], the strange quark contribution to the anomalous magnetic moment and the electromagnetic radius [3, 4, 5] and the scalar strange quark density of the nucleon [2, 3, 6]. In this note we will concentrate on the latter one.

A low energy theorem derived from current algebra relates the empirical isospin even pion-nucleon scattering amplitude at the Cheng-Dashen point to the πN sigma term [7]. By analysing the currently available experimental data, taking into account the strong

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t-dependence of the scalar form factor $\sigma(t)$ (with nucleon spinors $u(P)$),

$$\sigma(t)\bar{u}(P')u(P) = m_0\langle P'|\bar{u}u + \bar{d}d|P\rangle, \quad t = (P - P')^2, \tag{1}$$

the authors of ref. [8] conclude:

$$\sigma \equiv \sigma(0) = (45 \pm 8)\ \mathrm{MeV}. \tag{2}$$

Here $m_0 = (m_u + m_d)/2$ is the isospin averaged current quark mass. If the nucleon is free of strange quarks, σ should be equal to the matrix element of the $SU(3)$-octet scalar quark density

$$\sigma_8 = m_0\langle P|\bar{u}u + \bar{d}d - 2\bar{s}s|P\rangle. \tag{3}$$

A straightforward calculation of baryon masses to first order in the difference of the current quark masses yields $\sigma_8 \approx 25$ MeV, while corrections from $SU(3)$-breaking terms of next order indicate a shift to larger values [9], $\sigma_8 = 35 \pm 5$ MeV. Writing $\sigma = \sigma_8/(1-y)$ one derives for the ratio y of strange to non-strange pairs in the nucleon:

$$y \equiv \frac{2\langle P|\bar{s}s|P\rangle}{\langle P|\bar{u}u + \bar{d}d|P\rangle} = 0.2 \pm 0.2. \tag{4}$$

This result would indicate a significant violation of the Zweig rule, e.g. of the intuitive assumption that the nucleon is free of strange quarks, although with a large error.

Several calculations of such strange quark matrix elements have been carried out in chiral quark models, e.g. the Nambu & Jona-Lasinio (NJL) model [10, 11, 12]. In these models the nucleon is composed of three quasi-particles, the constituent quarks, which have a non-trivial structure governed by the mechanism of spontaneous chiral symmetry breaking. Since the nucleon consists of constituent u- and d-quarks, isospin symmetry implies that one can replace the nucleon state $|P\rangle$ in (4) by a constituent u-quark state $|U\rangle$, i.e. $y = 2\langle U|\bar{s}s|U\rangle / \langle U|\bar{u}u + \bar{d}d|U\rangle$. We will first outline problems in the calculation of scalar strangeness matrix elements in such models due to their sensitivity to the parametrization

and the regularization procedure. Next we study an extreme instanton scenario. Then we investigate the role of perturbative meson loop corrections in building up strange quark pairs in the nucleon.

a) NJL approach.

In the first part we consider the 3-flavor version of the NJL model [12]. It starts from the effective Lagrangian

$$\mathcal{L}_{NJL} = \bar{\psi}(i\partial^{\mu}\gamma_{\mu} - \hat{m})\psi + \mathcal{L}_4 + \mathcal{L}_6, \tag{5}$$

with $\psi = (u, d, s)^t$ and the current quark mass matrix $\hat{m}$ =diag(m_u, m_d, m_s). The local four-quark interaction $\mathcal{L}_4$ is symmetric under the chiral $U(3)_L \times U(3)_R$ group:

$$\mathcal{L}_4 \;\; = \;\; G_S \left[(\bar{\psi}\frac{\lambda_a}{2}\psi)^2 + (\bar{\psi}i\gamma_5\frac{\lambda_a}{2}\psi)^2 \right] + ..., \tag{6}$$

where additional terms need not be considered in the present context since they do not enter in the expressions for scalar densities at the mean field level. Here λ_a with $a = 0, 1, ..., 8$ are the standard $U(3)$ flavor matrices including the singlet $\lambda_o = \sqrt{2/3}$ diag$(1, 1, 1)$.

In nature the axial $U(1)_A$ symmetry is broken dynamically, presumably by instantons. A minimal effective interaction, suggested by 't Hooft [13], which selectively breaks $U(1)_A$ but leaves the remaining $SU(3)_L \times SU(3)_R \times U(1)_V$ untouched, is a 6-quark interaction in the form of a flavor-mixing 3×3 determinant:

$$\mathcal{L}_6 = G_D \left\{ \det[\bar{\psi}_i(1+\gamma_5)\psi_j] + \det[\bar{\psi}_i(1-\gamma_5)\psi_j] \right\}. \tag{7}$$

The effective Lagrangian (5) with the interaction $\mathcal{L}_4 + \mathcal{L}_6$ has been used extensively in the mean field approximation to study a variety of low energy, non-perturbative phenomena. A cutoff Λ of order 1 GeV is employed to regularize momentum space (loop) integrals. The physical picture behind this model is that strong interactions between quarks operate

at low momenta, i.e. for quark momenta smaller than Λ, whereas they are "turned off" for momenta larger than Λ.

For sufficiently strong coupling the vacuum undergoes spontaneous chiral symmetry breaking (SCSB). Quark condensates $\langle \bar{u}u \rangle$, $\langle \bar{d}d \rangle$ and $\langle \bar{s}s \rangle$ develop. Current quarks turn into constituent quarks with large dynamical masses determined by a set of gap equations. For example,

$$M_u = m_u - G_S \langle \bar{u}u \rangle - G_D \langle \bar{d}d \rangle \langle \bar{s}s \rangle, \tag{8}$$

with the quark condensates ($q = u, d, s$; $N_c = 3$):

$$\langle \bar{q}q \rangle = -4N_c i \int^{\Lambda} \frac{\mathrm{d}^4 p}{(2\pi)^4} \frac{M_q}{p^2 - M_q^2 + i\epsilon}. \tag{9}$$

The use of a four-momentum cutoff Λ_4 gives

$$\langle \bar{q}q \rangle = -\frac{N_c}{4\pi^2} M_q \left[\Lambda_4^2 - M_q^2 \ln(1 + \frac{\Lambda_4^2}{M_q^2}) \right], \tag{10}$$

while employing a three-momentum cutoff Λ_3 yields

$$\langle \bar{q}q \rangle = -\frac{N_c}{2\pi^2} M_q \left[\Lambda_3 \sqrt{\Lambda_3^2 + M_q^2} - M_q^2 \,\mathrm{arsinh}(\Lambda_3/M_q) \right]. \tag{11}$$

Using cutoffs in the range $(0.6 - 0.9)$ GeV and constituent quark masses $M_q = (0.3 - 0.4)$ GeV, the values for the quark condensates come out typically as $\langle \bar{q}q \rangle \approx -(250\mathrm{MeV})^3$.

SCSB implies the existence of pseudoscalar Goldstone bosons. In the chiral limit $m_u = m_d = m_s = 0$ and with $\mathcal{L}_6 = 0$ the whole pseudoscalar nonet of π's, K's, η_o and η_8 is massless. In the NJL model these modes emerge as explicit solutions of the quark-antiquark Bethe-Salpeter equation [12]. Dynamical $U(1)_A$ symmetry breaking by $\mathcal{L}_6$ of eqn.(7) gives the singlet η_o a non-zero mass. Furthermore, explicit breaking of chiral $SU(3)_L \times SU(3)_R$ symmetry by bare quark masses $m_s > m_{u,d} > 0$ moves all masses of the pseudoscalar nonet to their physical values. The $\eta - \eta'$ system is reproduced including its

mixing angle $\theta \approx -10°$ (26). In the "standard" NJL model the constituent u-quark mass remains as a free parameter which we can choose as $M_u \approx 330$ MeV, about one third of the nucleon mass.

The constituent quarks are quasi-particles. Their strong interaction dresses the valence quarks by quark-antiquark polarization clouds, so that the constituent quarks have non-trivial formfactors. In particular, the matrix element of the scalar quark density of flavor q in a constitutent u-quark is given by the Feynman-Hellmann theorem [14],

$$\langle U|\bar{q}q|U\rangle = \frac{\partial M_u}{\partial m_q}. \tag{12}$$

With eq. (8) we obtain the following explicit expression for the $\bar{s}s$ content of the u-quark quasi-particle:

$$\langle U|\bar{s}s|U\rangle = -G_S\frac{\partial\langle\bar{u}u\rangle}{\partial m_s} - G_D\left[\langle\bar{s}s\rangle\frac{\partial\langle\bar{d}d\rangle}{\partial m_s} + \langle\bar{d}d\rangle\frac{\partial\langle\bar{s}s\rangle}{\partial m_s}\right]. \tag{13}$$

Note that the important term is the last one proportional to $\partial\langle\bar{s}s\rangle/\partial m_s$. Its coefficient $G_D < \bar{d}d >$ is well constrained by the physics of the $\eta - \eta'$-system. However, the value of the derivative

$$\frac{\partial\langle\bar{s}s\rangle}{\partial m_s} = \frac{\partial\langle\bar{s}s\rangle}{\partial M_s}\frac{\partial M_s}{\partial m_s} \tag{14}$$

depends strongly on the regularization procedure. Using expression (10) with a four-momentum cutoff one finds that $\langle\bar{q}q\rangle$ exhibits a minimum at $M_q/\Lambda_4 \approx 0.75$. Hence with typical cutoffs $\Lambda_4 \approx 0.8$ GeV and $M_s \approx (0.5 - 0.6)$ GeV, the derivative (14) is extremely small, and the resulting ratio of strange to non-strange pairs turns out to be $y < 0.02$ in this case, with a strong sensitivity to the parameter M_s/Λ_4 [11]. In contrast the three-momentum cutoff [10] generally leads to much larger scalar strange pair admixtures than the four-momentum cutoff [12]. In table I we show typical results obtained with a three-momentum cutoff.

	G_S (GeV^{-2})	$G_D\langle\bar{s}s\rangle$ (GeV^{-2})	M_u (GeV)	M_s (GeV)	Λ (GeV)
NJL	16.4	4.1	0.33	0.52	0.65
INS	0	36.3	0.54	0.64	0.57

	$\langle U\|\bar{u}u\|U\rangle$	$\langle U\|\bar{d}d\|U\rangle$	$\langle U\|\bar{s}s\|U\rangle$	y
NJL	2.2	0.6	0.2	0.15
INS	1.4	0.6	0.6	0.6

TABLE I: *Upper part:* input for the NJL and the instanton (INS) model with three-momentum cutoff. The parameters are adjusted to reproduce quark condensates and the pseudoscalar meson spectrum including decay constants. *Lower part:* resulting scalar density matrix elements for the constituent u-quark and the strange pair fraction $y = 2\langle U|\bar{s}s|U\rangle/\langle U|\bar{u}u + \bar{d}d|U\rangle$.

b) Instanton approach.

Some years ago Diakonov and Petrov [15] suggested that SCSB in QCD may be entirely generated by instantons. If so, the 't Hooft interaction (7) should dominate in the low energy regime of QCD. Motivated by their work we have studied a version of the NJL model in which $\mathcal{L}_4$ vanishes ($G_S = 0$) and SCSB is generated only by the flavor mixing term proportional to G_D in (8). A careful search for minima of the effective potential[2] in the three-momentum regularization scheme shows that this particular pattern of SCSB requires a coupling strength which comes out quite stable around $G_D \approx 140 \cdot \Lambda^{-5}$. After fixing the pion decay constant to its physical value $f_\pi = 93$ MeV by adjusting the cutoff at $\Lambda \approx 0.57$ GeV, the constituent quark masses turn out to be rather large, namely $M_u \approx 0.5$ GeV and $M_s \approx 0.6$ GeV. Note that the cutoff for the 't Hooft interaction should be related to the average instanton size ρ (in fact it should be compared with $1/\rho \approx 0.6$ GeV from [15]). Current quark masses and quark condensates are in quite good agreement with standard values once the empirical pion and kaon masses are reproduced. The η-meson mass $m_\eta \approx 0.57$ GeV comes out close to its empirical value. However, in the η'-channel the t'Hooft interaction is repulsive and cannot generate a bound state.

We now present numerical results for the NJL and the pure instanton scenario (INS) in comparison (see table I). The enhanced flavor mixing of the INS reflects the much larger strength of the 't Hooft interaction (7) as compared to the one in the "standard" NJL approach. Hence, if chiral flavor dynamics is dominated by instantons, one confronts ratios y as large as $1/2$. With more conventional versions of the NJL model we find ratios $y \lesssim 0.15$. However, the strong dependence on the regularization scheme prohibits more

[2]The effective potential of the NJL model is worked out in detail in [16] and reads for $G_S = 0$: $\epsilon(M_u, M_d, M_s) = \sum_{q=u,d,s} w_q^0 - G_D \langle \bar{u}u \rangle \langle \bar{d}d \rangle \langle \bar{s}s \rangle$, with its free part: $w_q^0 = -(M_q - m_q)\langle \bar{q}q \rangle + N_c i \mathrm{tr} \int \frac{d^4p}{(2\pi)^4} \ln(-i \not{p} - M_q + i\epsilon)$.

FIGURE.1: Self-energy of a u-quark in the presence of a kaon cloud.

reliable estimates, so that these numbers should be considered as upper limits.

c) Kaon loops.

Our next task is to consider perturbative corrections to the quark propagator due to the emission and reabsorption of pseudoscalar mesons. The admixture of strange pairs in such processes comes from dressing the quark with a kaon cloud as shown in figure 1. We first convince ourselves that the shift δM_u of the u-quark mass from such meson loops is small so that perturbation theory is justified. Then we use eq. (12) again and calculate the correction

$$\delta\langle U|\bar{s}s|U\rangle = \frac{\partial(\delta M_u)}{\partial m_s} \tag{15}$$

to the strange quark admixture in the constituent u-quark from such mechanisms. Perturbative corrections to constituent quark masses have also been investigated in [17] within the framework of a model restricted to $SU(2)$. Furhermore Koepf et al. [4] used the $SU(3)$-cloudy bag model for determining the strange magnetic form factor and the strange axial charge of the nucleon. A recent work [18] employs light-cone meson-nucleon vertex functions for calculating axial form factors of the nucleon following similar ideas.

In our calculations we use the semi-bosonized version of the $SU(3)$ NJL model (5,6) with $\bar{\psi}i\gamma_5(\lambda_a/2)\psi$ replaced by the corresponding pseudoscalar mesons treated as collective degrees of freedom. The self-energy of a u-quark of momentum p dressed by a kaon cloud

as in figure 1 is

$$\Sigma(p) = ig^2 \int \frac{d^4q}{(2\pi)^4} i\gamma_5 \frac{\not{q}_+ + M_s}{q_+^2 - M_s^2 + i\epsilon} i\gamma_5 \frac{1}{q_-^2 - m_K + i\epsilon} \equiv M_s A(p^2) - \not{p} B(p^2), \tag{16}$$

where q is the loop momentum and $q_\pm = q \pm p/2$; M_s is the mass of the strange constituent quark and m_K the kaon mass in the intermediate state. The kaon-quark coupling constant g is given to leading order in the pseudoscalar meson mass by the Goldberger-Treiman relation:

$$g = \frac{M_u + M_s}{2f_K} + \mathcal{O}(m_K^2), \tag{17}$$

where f_K is the kaon decay constant. The latter is related to the quark condensates and the current quark masses (in leading order) by the Gell -Mann, Oakes, Renner (GOR) relation,

$$(f_K m_K)^2 = -\frac{1}{2}(m_u + m_s)\left(\langle \bar{u}u \rangle + \langle \bar{s}s \rangle\right) + \mathcal{O}(m_K^4). \tag{18}$$

The scalar functions $A(p^2)$ and $B(p^2)$ are given in terms of loop integrals

$$I_1(\mu) \equiv 2i \int \frac{d^4q}{(2\pi)^4} \frac{1}{q^2 - \mu^2 + i\epsilon} \tag{19}$$

and

$$I_2(\mu, \mu'; p^2) \equiv i \int \frac{d^4q}{(2\pi)^4} \frac{1}{q_+^2 - \mu^2 + i\epsilon} \cdot \frac{1}{q_-^2 - \mu'^2 + i\epsilon}, \tag{20}$$

as follows:

$$A(p^2) = -g^2 I_2(M_s, m_K; p^2), \tag{21}$$

$$B(p^2) = -\tfrac{1}{2} g^2 \left[\left(1 + \tfrac{M_s^2 - m_K^2}{p^2}\right) I_2(M_s, m_K; p^2) + \tfrac{1}{2p^2}\left(I_1(m_K) - I_1(M_s)\right) \right].$$

Explicit expressions for $I_{1,2}$ are given in ref. [12]. For consistency we employ the same regularization procedure as used in the primary NJL model which generates the constituent quarks and the pseudoscalar mesons entering in (16). The perturbative correction from

(16) to the constituent u-quark mass becomes:

$$\delta M_u = M_s A(M_u^2) - M_u B(M_u^2). \tag{22}$$

Numerically one finds $A(M_u^2) \approx 0.2$ and $A \approx 2B$ so that the relative correction $\delta M_u/M_u$ of the u-quark mass due to the K^+ cloud is about 10%. This justifies the use of perturbation theory.

In the absence of $U(1)_A$ breaking effects ($G_D = 0$ in the gap equation (8)), the process $u \to (u\bar{s})s$ which turns the u-quark into a s-quark and a K^+-meson is the only one that contributes to $\langle U|\bar{s}s|U\rangle$. With $G_D \neq 0$ there are additional non-leading corrections involving pion cloud contributions to the quark self-energy as well, through $\bar{s}s$-dressings of their constituent quarks, but they turn out to be negligibly small.

In our estimate of K^+ loop effects we therefore use $G_D = 0$ and adjust G_S in (6), the current quark masses and the cutoff in such a way that $M_u = 330$ MeV and the pion mass, the kaon mass and kaon decay constant $f_K = 114$ MeV coincide with their empirical values. Following eq. (12) we then calculate

$$\delta\langle U|\bar{s}s|U\rangle = \frac{\partial}{\partial m_s}\left[M_s A(M_u^2)\right] - \frac{\partial}{\partial m_s}\left[M_u B(M_u^2)\right]. \tag{23}$$

Several effects cooperate in this expression. First the constituent quark mass M_s grows with the current quark mass m_s according to the gap equation (8). Secondly, the K^+ mass is related to m_s by eq. (18) and the resulting contribution to (23) is negative. Furthermore, the Goldberger-Treiman relation (17) together with the gap equation implies a positive derivative of the kaon-quark coupling constant with respect to m_s. Altogether, the first effect prevails, and we end up with

$$\delta\langle U|\bar{s}s|U\rangle \approx 0.03, \tag{24}$$

using a three-momentum cutoff scheme. About half of this value results when a four-momentum cutoff is used.

Diagonal matrix elements such as $\delta\langle U|\bar{u}u|U\rangle$ from pion and kaon loops turn ot to be of the same order of magnitude as (24). Hence $\langle U|\bar{u}u+\bar{d}d|U\rangle$ is not substantially modified from its mean field value in table I, calculated by varying the gap equation (8) with respect to $m_{u,d}$. Hence the correction from the K^+ cloud to the ratio y becomes:

$$\delta y \approx \frac{2\delta\langle U|\bar{s}s|U\rangle}{\langle U|\bar{u}u+\bar{d}d|U\rangle} \lesssim 0.03. \tag{25}$$

d) <u>η and η' loops.</u>

Loops involving the η and η' mesons contribute to δy as well, but their effects are small, as we will now show. First, the contribution of the η' is suppressed due to its large mass. However, the η mass is not much larger than the kaon mass, so that the corresponding loop correction has to be investigated more carefully.

The η has the following decomposition:

$$\eta = \eta_u\left(\bar{u}u+\bar{d}d\right) - \eta_s\bar{s}s, \tag{26}$$

where $\eta_u = \frac{1}{\sqrt{3}}\left(\frac{\cos\theta}{\sqrt{2}} - \sin\theta\right)$, $\eta_s = \frac{1}{\sqrt{3}}\left(\sqrt{2}\cos\theta + \sin\theta\right)$ in terms of the $\eta-\eta'$ mixing angle θ. For $\theta = -10^\circ$ we have $\eta_u \approx 0.5$, $\eta_s \approx 0.7$ Note the reduction by a factor η_u of the η coupling to a u-quark as compared to that of a kaon. Altogether it turns out that the correction $(\delta M_u)_\eta$ to the constituent u-quark mass due to the η cloud is only about 5 MeV (compared to 35 MeV for the kaon cloud). Next, consider

$$\delta\langle U|\bar{s}s|U\rangle_\eta = \frac{\partial}{\partial m_s}(\delta M_u)_\eta = \frac{\partial(\delta M_u)_\eta}{\partial m_\eta}\frac{\partial m_\eta}{\partial m_s} + \frac{\partial(\delta M_u)_\eta}{\partial \eta_u}\frac{\partial \eta_u}{\partial m_s} = \frac{\partial(\delta M_u)_\eta}{\partial m_\eta}\frac{\partial m_\eta}{\partial m_s} + \frac{2\cdot(\delta M_u)_\eta}{\eta_u}\frac{\partial \eta_u}{\partial m_s}. \tag{27}$$

The first term on the r.h.s. involves the positve derivative of the η mass m_η with respect to the strange quark mass together with the (negative) change of the η loop integral when changing m_η. This product has a small numerical value (about -0.01). The last term (in which the relation $(\delta M_u)_\eta \propto \eta_u^2$ has been used) reflects the dependence of the $\eta-\eta'$

mixing pattern on the strange quark mass. For $m_s = m_u = 0$ we have $\theta = 0$. Assuming a linear dependence $\theta \propto m_s$ we estimate at $\theta \approx -10°$ with $m_s \approx 130$ MeV:

$$\frac{2(\delta M_u)_\eta}{\eta_u}\frac{\partial \eta_u}{\partial m_s} \approx -10\,\mathrm{MeV}\cdot\frac{\sin\theta+\sqrt{2}\cos\theta}{\cos\theta-\sqrt{2}\sin\theta}\cdot\frac{\partial\theta}{\partial m_s} \approx 0.013. \tag{28}$$

Adding up both terms in eq. (27) gives a negligibly small contribution to $\langle U|\bar{s}s|U\rangle$, about one order of magnitude smaller than that from the kaon cloud.

In summary we have analysed strange quark admixtures to the scalar density of the nucleon in terms of $\bar{s}s$-components in the quasi-particle structure of constituent quarks. In our approach such components arise from the non-perturbative dressing of the quarks by scalar mean fields, and from perturbative kaon cloud effects. We find that in the "standard" NJL model with moderate axial $U(1)$ breaking, the upper limit for the ratio y of $\bar{s}s$ pair admixtures in the nucleon relative to $\bar{u}u$ and $\bar{d}d$ pairs is about 0.15 (with substantial uncertainties due to the strong dependence on details of the regularization procedure). In contrast, if instantons dominate the low energy dynamics so that the effective interaction is governed by axial $U(1)$ breaking, the resulting $\bar{s}s$ admixtures can be much larger, but with similar uncertainties. Kaon cloud effects alone, on the other hand, would not give large strange pair components in the nucleon. We find an upper limit of about 3% from this source.

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